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# Light-cone approach to eleven dimensional supergravity[1]

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**Abstract**

Manifestly supersymmetric formulation of eleven dimensional supergravity in the framework of light-cone approach is discussed.

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# 1 Motivation and summary of results

The long term motivation for our study of $11d$ supergravity is related to conjectured iterrelation between superstring theory and AdS higher spin massless field theory. Ten years ago E.S. Fradkin [1], based on studies made in [2]-[5], put forward the idea that string theory and anti-de Sitter higher spin gauge theory, though different, eventually may turn out to be different phases of one and the same unified field theory with new forces mediated by higher spin gauge fields . According to this conjecture string theory can be interpreteted as resulting from some kind of spontaneous breakdown of higher spin symmetries.

To develop this idea it was conjectured recently [6] that *superstrings could be considered as living at the boundary of* $11$*-dimensional AdS space while their unbroken (symmetric) phase is realized as a theory of higher spin massless fields living in this* $AdS_{11}$ *space*. Some discussion of this theme can be found in [7] where it was demonstrated that if one restricts attention to totally symmetric fields and make some mild assumption about the (spontaneously) broken form of AdS theory Hamiltonian then *leading components of AdS massless totally symmetric arbitrary spin states become massive string states belonging to leading Regge trajectory.*

As is well known the standard $11d$ supergravity [8] does not admit an extension with a cosmological constant, i.e. does not have $AdS_{11}$ vacuum [9](see also [10, 11]). One other hand, in [12] certain massless $AdS_{11}$ graviton supermultiplet was found[2]. This novel supermultiplet contains fields of the usual $11d$ supergravity plus additional ones. One can expect that these additional fields may allow one to overcome no-go theorem and construct a consistent supergravity admitting $AdS_{11}$ ground state[3].

The first step in this direction would be to find free field theoretic realization of this $AdS_{11}$ supermultiplet and then try to construct interactions. Light-cone approach provides self-contained setup to study these questions. The main advantage of light-cone approach is that it allows one to discuss supersymmetric theories in terms of unconstrained scalar superfields. Before attempting to study $AdS_{11}$ supergravity it would be interesting to consider the usual $11d$ supergravity which, to our knowledge was not previously discussed in superfield light-cone gauge.

To discuss $11d$ supergravity we will exploit the method of [16] which reduces the problem of finding a new (light-cone gauge) dynamical system to the problem of finding a new solution of commutation relations of the defining symmetry algebra (in our case $11d$ Poincaré superalgebra). In the past this method was successfully applied for finding manifestly supersymmetric formulations of various theories [17,

[2]Related interesting discussion can be found in [13].

[3] Certain massless $AdS_{11}$ graviton multiplet is also predicted by eleven dimensional version of $AdS_{10}$ higher spin gauge theories discovered in [14]. These theories allow more or less straightforward generalization to $AdS_{11}$ [15]. Since normally a tower of infinite higher spin fields contains of supergravity multiplet one expects that eleven dimensional version of theories discussed in [14] also describes some $AdS_{11}$ graviton multiplet.

18, 19][4]. Despite many known examples of cubic vertices given in the literature, constructing cubic vertices for concrete field theories is still a challenging procedure. A general method essentially simplifying the procedure of obtaining cubic interaction vertices was discovered in [24], developed in [25, 26] and formulated finally in [27]. One of the characteristic features of this method is reducing manifest transverse $so(d-2)$ invariance (which is $so(9)$ for $11d$ supergravity) to $so(d-4)$ invariance (which is $so(7)$ in this paper)[5]. On the other hand, it is $so(7)$ symmetry that is manifest symmetry of unconstrained superfield formulation of $11d$ supergravity. In other words the manifest symmetries of our method and the one of unconstrained superfield formulation of $11d$ supergravity match. Here we demonstrate how the method of Ref. [27] works for the case of $11d$ supergravity.

Light-cone gauge $11d$ supergravity can be formulated in light-cone superspace which is based on position coordinates $x^\mu$, and Grassmann position coordinates $\theta^\alpha$ [6]. In this light-cone superspace we introduce scalar superfield $\Phi(x^\mu,\theta)$. Instead of position space it is convenient to use momentum space for all coordinates except the light-cone time $x^+$. This implies using $p^+$, $p^R$, $p^L$, $p^i$, $\lambda^\alpha$, instead of $x^-$, $x^L$, $x^R$, $x^i$, $\theta^\alpha$ respectively. Thus we consider the scalar superfield $\Phi(x^+,p^+,p^R,p^L,p^i,\lambda)$ with the following expansion in powers of Grassmann momentum $\lambda$

$$\begin{aligned}\Phi(p,\lambda) &= \beta^2 A + \beta\lambda^\alpha\psi^\alpha + \beta\lambda^{\alpha_1}\lambda^{\alpha_2}A^{\alpha_1\alpha_2}\\ &+ \lambda^{\alpha_1}\lambda^{\alpha_2}\lambda^{\alpha_3}\psi^{\alpha_1\alpha_2\alpha_3} + \lambda^{\alpha_1}\dots\lambda^{\alpha_4}A^{\alpha_1\dots\alpha_4} + \frac{1}{\beta}(\epsilon\lambda^5)^{\alpha_1\alpha_2\alpha_3}\psi^{\alpha_1\alpha_2\alpha_3 *}\\ &- \frac{1}{\beta}(\epsilon\lambda^6)^{\alpha_1\alpha_2}A^{\alpha_1\alpha_2 *} - \frac{1}{\beta^2}(\epsilon\lambda^7)^\alpha\psi^{\alpha *} + \frac{1}{\beta^2}(\epsilon\lambda^8)A^*\,, \end{aligned}\tag{1}$$

where we use the notation[7]

$$(\epsilon\lambda^{8-n})^{\alpha_1\dots\alpha_n} \equiv \frac{1}{(8-n)!}\epsilon^{\alpha_1\dots\alpha_n\alpha_{n+1}\dots\alpha_8}\lambda^{\alpha_{n+1}}\dots\lambda^{\alpha_8}\tag{2}$$

[4]Derivation of light-cone formulation from covariant Largangians of $N=4$ SYM was discussed in [20, 21]. Discussion of higher spin massless fields may be found in [22]-[27].

[5] Previously, reducing the manifest $so(d-2)$ symmetry to $so(d-4)$ was used to formulate superfield theory of $IIA$ superstrings [19]. There this reducing was motivated by the desire to get unconstrained superfield formulation. In [27] the main motivation for reducing was the desire to get the most general solution for cubic vertex for arbitrary spin fields of (super) Poincaré invariant theory. Discussion of $so(7)$ formalism in the context of M(atrix) theory can be found in [28].

[6]$\mu = 0,1,\dots 10$ are $so(10,1)$ vector indices, $\alpha=1,\dots,8$ are $so(7)$ spinor index, $I,J,K=1,\dots 9$ are $so(9)$ transverse indices, $i,j,k=1,\dots,7$ are $so(7)$ transverse indices. Coordinates in light-cone directions are defined by $x^\pm \equiv (x^{10}\pm x^0)/\sqrt{2}$. Remaining transverse coordinates $x^I$ are decomposed into $x^i$, $x^{R,L}$ where $x^{R,L}\equiv(x^8\pm \mathrm{i}x^9)/\sqrt{2}$. The scalar product of two $so(9)$ vectors is decomposed then as $X^IY^I = X^iY^i + X^RY^L + X^LY^R$. For momentum in light-cone direction we use simplified notation $\beta\equiv p^+$.

[7] In what follows a momentum $p$ as argument of the superfield $\Phi$ and $\delta$–functions designates the set $\{p^I,\beta\}$. Also we do not show explicitly the dependence of the superfield on evolution parameter $x^+$.

and $\epsilon^{\alpha_1\dots\alpha_8}$ is the Levi-Civita tensor. The only constraint which the superfield $\Phi$ should satisfy is the reality constraint

$$\Phi(-p,\lambda) = \beta^4 \int d^8\lambda^\dagger e^{\lambda\lambda^\dagger/\beta} (\Phi(p,\lambda))^\dagger \,. \tag{3}$$

This constraint tells us that some fields in (1) are related by Hermitean conjugation. In (1) the component fields carrying even number of spinor indices describe bosonic fields

$$A^{\alpha_1\dots\alpha_4}(70) \sim \{h^{ij}(27^0), h^{RL}(1^0), C^{ijk}(35^0), C^{RLi}(7^0)\}\,, \tag{4}$$

$$A^{\alpha_1\alpha_2}(28) \sim \{h^{Li}(7^{-1}), C^{Lij}(21^{-1})\}\,, \qquad A = h^{LL}/\sqrt{2}\,, \tag{5}$$

while the fields with odd number of spinor indices are responsible for gravitino field. Superscripts in (4),(5) indicate $J^{RL}$ charge. Light-cone gauge action for the both free and interacting theory takes then the following standard form

$$S = \int dx^+ \beta d\beta d^9 p d^8\lambda \Phi(-p,-\lambda) \mathrm{i}\partial^- \Phi(p,\lambda) + \int dx^+ P^- \,, \tag{6}$$

where $\partial^- = \partial/\partial x^+$ and $P^-$ is Hamiltonian. For free theory $P^-$ is given by the standard expression

$$P^-_{(2)} = \int \beta d\beta d^9 p d^8\lambda \Phi(-p,-\lambda) \left(-\frac{p^I p^I}{2\beta}\right) \Phi(p,\lambda)\,. \tag{7}$$

Now let us discuss cubic interactions. General structure of 3 point interaction vertices is obtainable from commutation relations of Poincaré superalgebra. Some of the latter lead to the following expression for the Hamiltonian

$$P^-_{(3)} = \int d\Gamma_3 \prod_{a=1}^3 \Phi(p_a,\lambda_a)\, p^-_{(3)}\,, \tag{8}$$

where the indices $a,b,c = 1,2,3$ label three interacting superfields and

$$d\Gamma_3 \equiv \delta^{10}(\sum_{a=1}^3 p_a)\delta^8(\sum_{a=1}^3 \lambda_a) \prod_{a=1}^3 d\beta_a d^9 p_a d^8\lambda_a \,. \tag{9}$$

The Hamiltonian density $p^-_{(3)}$ depends on momenta $\beta_a$, transverse momenta $p^I_a$ and Grassmann momenta $\lambda_a$. The $\delta$- functions in (9) respect conservation laws for these momenta.

Next, using commutation relations of Hamiltonian with $J^{+I}$ and certain supercharges one finds that the Hamiltonian density $p^-_{(3)}$, depends on momenta $p^I_a$ and $\lambda_a$ in a special manner. Namely, it turns out that $p^-_{(3)}$ depends on $p^I_a$ and $\lambda_a$ through the following quantities

$$\mathbf{P}^I_{ab} \equiv p^I_a \beta_b - p^I_b \beta_a\,, \qquad \Lambda_{ab} \equiv \lambda_a \beta_b - \lambda_b \beta_a\,. \tag{10}$$

The remarkable simplification is that the new momenta $\mathbf{P}^I_{12}$, $\mathbf{P}^I_{23}$, $\mathbf{P}^I_{31}$ are not independent: all of them are expressible through $\mathbf{P}^I$ defined by[8]

$$\mathbf{P}^I = \frac{1}{3}\sum_{a=1}^{3} \check{\beta}_a p^I_a\,, \qquad \check{\beta}_a \equiv \beta_{a+1} - \beta_{a+2}\,, \quad \beta_a \equiv \beta_{a+3}\,. \tag{11}$$

The same happens for Grassmann momenta, i.e. due to momentum conservation laws for $\beta_a$ and Grassmann momentum $\lambda_a$ the new Grassmann momenta $\Lambda_{12}$, $\Lambda_{23}$, $\Lambda_{31}$ (see 10) are expressible in terms of one momentum $\Lambda$ defined by

$$\Lambda = \frac{1}{3}\sum_{a=1}^{3} \check{\beta}_a \lambda_a\,. \tag{12}$$

The usage of $\mathbf{P}^I$ and $\Lambda$ is advantageous since they are invariant under cyclic permutation of indices $1, 2, 3$ which label three interacting fields. Thus $p^-_{(3)}$ is eventually the function of $\mathbf{P}^I$, $\Lambda$ and $\beta_a$,

$$p^-_{(3)} = p^-_{(3)}(\mathbf{P}, \Lambda, \beta_a)\,. \tag{13}$$

The $p^-_{(3)}$, by definition, is a monomial of degree $k$ in $\mathbf{P}^I$. As is well known the original $11d$ supergravity is described by the vertex $p^-_{(3)}$ involving terms of second order in transverse momentum $\mathbf{P}^I$, i.e. we have to set $k = 2$ . Let us for flexibility however keep $k$ to be arbitrary. Then the cubic vertex can be presented as

$$p^-_{(3)} = \mathbf{P}^{I_1} \ldots \mathbf{P}^{I_k} p^{-I_1 \ldots I_k}_{(3)}(\Lambda, \beta_a)\,. \tag{14}$$

In general, the $p^{-I_1\ldots I_k}_{(3)}$ is a complicated $so(9)$ tensor depending on Grassmann momentum $\Lambda$ and light-cone momenta $\beta_a$. It is the finding this tensor that is the most difficult part of analysis of cubic vertices. Note that $p^-_{(3)}$ after reducing to $so(7)$ notation has the decomposition

$$p^-_{(3)} = (\mathbf{P}^L)^k p^{-R\ldots R}_{(3)} + (\mathbf{P}^L)^{k-1}\mathbf{P}^i p^{-iR\ldots R}_{(3)} + \ldots + (\mathbf{P}^R)^k p^{-L\ldots L}_{(3)}\,. \tag{15}$$

In [27] a method was suggested which allows one to express $p^{-I_1\ldots I_k}_{(3)}$, which is $so(9)$ tensor, in terms of vertex $\tilde{V}_0$, which is $so(7)$ scalar and has charge $k$ with respect to $J^{RL}$. The general formula is

$$p^-_{(3)}(\mathbf{P}, \Lambda, \beta_a) = (\mathbf{P}^L)^k E_q E_\rho \tilde{V}_0(\Lambda, \beta_a)\,, \tag{16}$$

where the operators $E_q$, $E_\rho$ are defined by relations

$$E_q = \exp(-q^j \mathbf{M}^{Lj}_\Lambda)\,, \tag{17}$$

[8]By using momentum conservation laws for $p^I_a$ and $\beta_a$ it is easy to check that $\mathbf{P}^I_{12} = \mathbf{P}^I_{23} = \mathbf{P}^I_{31} = \mathbf{P}^I$.

$$E_\rho \equiv \sum_{n=0}^{k} (-\rho)^n \frac{\Gamma(\frac{7}{2}+k-n)}{2^n n! \Gamma(\frac{7}{2}+k)} (\mathbf{M}_\Lambda^{Lj}\mathbf{M}_\Lambda^{Lj})^n \,, \tag{18}$$

and we use the notation

$$q^i \equiv \frac{\mathbf{P}^i}{\mathbf{P}^L}\,, \qquad \rho \equiv \frac{\mathbf{P}^i\mathbf{P}^i + 2\mathbf{P}^R\mathbf{P}^L}{2(\mathbf{P}^L)^2}\,, \qquad \frac{\mathbf{P}^R}{\mathbf{P}^L} = \rho - \frac{q^2}{2}\,. \tag{19}$$

The vertex $\tilde{V}_0$ satisfies the following equations

$$(\mathbf{M}_\Lambda^{RL} - k)\tilde{V}_0 = 0\,, \qquad \mathbf{M}_\Lambda^{Ri}\tilde{V}_0 = 0\,, \qquad \mathbf{M}_\Lambda^{ij}\tilde{V}_0 = 0\,. \tag{20}$$

and it depends only on Grassmann momentum $\Lambda$ and light-cone momenta $\beta_a$. The dependence on the transverse space momentum $\mathbf{P}^I$ is thus isolated explicitly.

The representation for $p^-_{(3)}$ given in Eqs.(16)-(20) is universal and valid for arbitrary (super) Poincaré invariant theory. In order to get cubic vertices one needs (i) to find solutions to (20); (ii) to insert $\tilde{V}_0$ and appropriate spin parts of angular momentum $\mathbf{M}^{IJ}$ fixed by representation theory of super Poincaré algebra in (16). For the case under considerations the appropriate $\mathbf{M}^{IJ}$ are given by

$$\mathbf{M}_\Lambda^{RL} = \frac{1}{2}\theta_\Lambda\Lambda - 2\,, \quad \mathbf{M}_\Lambda^{Ri} = -\frac{1}{2\sqrt{2}}\hat{\beta}\theta_\Lambda\gamma^i\theta_\Lambda\,, \quad \mathbf{M}_\Lambda^{Li} = \frac{1}{2\sqrt{2}\hat{\beta}}\Lambda\gamma^i\Lambda\,, \tag{21}$$

where[9]

$$\hat{\beta} \equiv \beta_1\beta_2\beta_3\,, \tag{22}$$

and the $\theta_\Lambda$ is defined by (anti)commutation relation $\{\theta_\Lambda, \Lambda\} = 1$.

The remarkable property of the Eqs. (20) is that normally they are quite simple to analyse[10]. For instance, for the case of under consideration all what one needs is to analyse the first equation in (20). Indeed, making use of expression for $\mathbf{M}^{RL}$ given in (21) we find

$$\Lambda\theta_\Lambda\tilde{V}_0 = 2(2-k)\tilde{V}_0\,. \tag{23}$$

The operator $\Lambda\theta_\Lambda$ counts the degree of Grassmann momentum $\Lambda$ involved in $\tilde{V}_0$ which, by definition, cannot involve terms of negative power in $\Lambda$ i.e. eigenvalues of $\Lambda\theta_\Lambda$ must be non-negative. This implies that vertices with terms higher than second order in $\mathbf{P}^I$, i.e. when $k > 2$, are forbidden. Note that it is the terms with $k = 4$ and $k = 6$ that would correspond to supersymmetric extension of higher derivative terms like $R^2_{....}$ and $R^3_{....}$. Therefore, the fact that vertices with $k = 4$ and $k = 6$ are forbidden implies that terms of second and third order in Rieman tensor do not allow

[9] $\gamma^i$ are usual $so(7)$ $\gamma$-matrices: $\{\gamma^i, \gamma^j\} = 2\delta^{ij}$. All of them are taken to be antisymmetric and hermitean.

[10] General solution to these equations can be found in [27].

supersymmetric extension[11]. Thus the only value of $k$ allowed by Eq.(23) is $k = 2$, i.e. $\tilde{V}_0 = const$[12], and this leads to cubic vertex of the original $11d$ supergravity

$$p^-_{(3)}(\mathbf{P}, \Lambda, \beta_a) = \frac{\kappa}{3}(\mathbf{P}^L)^2 E_q E_\rho\,, \tag{24}$$

where $\kappa$ is the gravitational constant. We choose normalization so that the cubic action for graviton field obtainable from (6),(24) coincides with the one of the Einstein-Hilbert

$$S_{EH} = \frac{1}{2\kappa^2}\int \sqrt{g}R\,, \tag{25}$$

$R = R^{\mu\nu}{}_{\mu\nu}$, $R^\mu{}_{\nu\lambda\sigma} = \partial_\lambda\Gamma^\mu_{\nu\sigma} + \ldots$, where we use the following expansion for metric tensor $g_{\mu\nu} = \delta_{\mu\nu} + \sqrt{2}\kappa h_{\mu\nu}$ and light-cone gauge $h^{+\mu} = 0$.

Making use of (24) and the formula for $\mathbf{M}^{Li}_\Lambda$ given in (21) we can work out the explicit representation for cubic vertex in a rather straightforward way

$$\begin{aligned}
\frac{3}{\kappa}p^-_{(3)} \;&=\; \mathbf{P}^{L2} - \frac{\mathbf{P}^L}{2\sqrt{2}\hat{\beta}}\Lambda\,\not{\mathbf{P}}\Lambda + \frac{1}{16\hat{\beta}^2}(\Lambda\,\not{\mathbf{P}}\Lambda)^2 - \frac{\mathbf{P}_I^2}{9\cdot 16\hat{\beta}^2}(\Lambda\gamma^j\Lambda)^2 \\
&+\; \frac{\mathbf{P}^R}{9\cdot 16\sqrt{2}\hat{\beta}^3}\Lambda\,\not{\mathbf{P}}\Lambda(\Lambda\gamma^j\Lambda)^2 + \frac{\mathbf{P}^{R2}}{2^7\cdot 63\hat{\beta}^4}((\Lambda\gamma^i\Lambda)^2)^2\,,
\end{aligned} \tag{26}$$

where $\not{\mathbf{P}} \equiv \mathbf{P}^i\gamma^i$, $\mathbf{P}_I^2 \equiv \mathbf{P}^I\mathbf{P}^I$. Thus we have two equivalent representations for $11d$ supergravity cubic interaction vertex given by (24) and (26). The representation (26) being manifest in Grassmann momentum $\Lambda$ is not convenient, however, in calculations. In contrast, the representation (24) does not show explicitly the dependence on $\Lambda$. However the remarkable feature of representation (24) is that it is expressed entirely in terms of spin operator $\mathbf{M}^{Li}$ which has clear algebraic properties. For this reason it is the representation (24) that is the most convenient in calculations. As compared to (26), the representation (24) is universal. For instance the cubic vertices of $IIA$ SURGA and $N = 1$ ten-dimensional SYM have similar form.

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[11] One important thing to note is that we proved absence of above mentioned higher derivative terms by using only commutation relations between Hamiltonian $P^-$ and kinematical generators. Kinematical generators, by definition, are the generators of super Poincaré algebra which have zero or positive $J^{+-}$ charge. It is reasonable to think that kinematical generators do not receive quantum corrections. If this indeed would be the case then our result could be considered as light-cone proof of nonrenormalization of $R^2_{....}$ and $R^3_{....}$ terms in $11d$ supergravity. Note that we discuss theory with 32 supercharges. Study of $R^3_{....}$ terms in string theory effective actions can be found in [29].

[12] Note that Eq.(23) for $k = 2$ tells us that $\tilde{V}_0$ does not depend on $\Lambda$ but then it still depends on light-cone momenta $\beta_a$. The fact that $\tilde{V}_0$ does not depend on $\beta_a$ too can be proved by using the requirement that all (super)charge densities are polynomial in transverse momentum $\mathbf{P}^I$.

## 2 Conclusion

We have discussed the light-cone gauge formulation of usual $11d$ supergravity. The formulation is given entirely in terms of light-cone scalar superfield allowing us to treat all component fields on an equal footing. Because the formalism we presented is algebraic in nature it can be extended to AdS spacetime in a relative straightforward way. Comparison of this formalism with other approaches available in the literature leads us to the conclusion that this is a very efficient formalism.

The formulation presented here should have a number of interesting applications and generalizations, some of which are:

(i) generalization to $AdS_{11}$ spacetime and study of massless $AdS_{11}$ graviton supermultiplet found in [12].

(ii) application of manifestly supersymmetric light-cone formalism to the study of the various aspects of M-theory along the lines [30]–[33];

(iii) generalization to cubic vertices of type IIB supergravity in $AdS_5 \times S^5$ background [34] and then to strings in this background [35, 36, 37].

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